

Unusual Higgs or Supersymmetry from Natural Electroweak Symmetry Breaking^{*}

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Abstract

This review provides an elementary discussion of electroweak symmetry breaking in the minimal and the next-to-minimal supersymmetric models with the focus on the fine-tuning problem – the tension between natural electroweak symmetry breaking and the direct search limit on the Higgs boson mass. Two generic solutions of the fine-tuning problem are discussed in detail: models with unusual Higgs decays; and models with unusual pattern of soft supersymmetry breaking parameters.

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I. INTRODUCTION

Minimal Supersymmetric Standard Model (MSSM) is a promising candidate for describing physics above the electroweak (EW) scale. The three gauge couplings unify [1, 2, 3, 4] at the grand unified theory (GUT) scale $\sim 2 \times 10^{16}$ GeV within a few percent, and the hierarchy between the EW scale and the GUT scale is naturally stabilized by supersymmetry (SUSY) [5]. In addition, if we add soft-supersymmetry-breaking terms (SSBs) at the GUT scale we typically find that the mass squared of the Higgs doublet which couples to the top quark (H_u), is driven to negative values at the EW scale. This triggers electroweak symmetry breaking [6] and the EW scale is naturally understood from SUSY breaking scale. Furthermore, assuming R-parity, the lightest supersymmetric particle (LSP) is stable and it is a natural candidate for dark matter of the universe.

The real virtue of supersymmetry is that the above mentioned features do not require any specific relations between soft-supersymmetry-breaking parameters and the only strong requirement on SUSY breaking scenarios is that these terms are of order the EW scale. However generic SSBs near the EW scale generically predict too light Higgs mass which is ruled out by LEP limits. The exact value of the Higgs mass is not relevant for low energy physics, nothing crucially depends on it, and yet, in order to stay above LEP limits ($m_h \gtrsim 114.4$ GeV) [7] the SSBs have to be either considerably above the EW scale and related to each other (or to the μ term) in a non-trivial way. SSBs can no longer be just generic which leads to strong requirements on possible models for SUSY breaking should these provide a natural explanation for the scale at which the electroweak symmetry is broken.

In this review we discuss in detail the electroweak symmetry breaking in the minimal and the next-to-minimal supersymmetric models and focus on the fine-tuning problem – the tension between natural electroweak symmetry breaking and the direct search limit on the Higgs boson mass. We will assume that the low scale supersymmetry is the scenario responsible for EWSB, and more importantly, the EWSB happens in a natural way without necessity of fine tuning.¹ We will see that this assumption will lead us to consider models in which the Higgs boson decays in an unusual way or models with unusual pattern of soft

¹ Of course it is possible that idea of low scale supersymmetry is wrong and there is a different mechanism realized in nature that is responsible for EWSB. For a review of ideas, see *e.g.* Ref. [8].

supersymmetry breaking parameters.

II. ELECTROWEAK SYMMETRY BREAKING IN THE MSSM, THE HIGGS MASS AND THE FINE TUNING PROBLEM

Let us start with reviewing reasons we expect superpartners at the EW scale. Confronting these expectations with constraints imposed by limits on the Higgs mass will lead us to the fine-tuning problem.

A. Electroweak symmetry breaking in the MSSM

The mass of the Z boson (or the EW scale), determined by minimizing the Higgs potential, is related to the supersymmetric Higgs mass parameter μ and the soft-SUSY-breaking mass-squared parameters for H_u and H_d by:

$$\frac{1}{2}M_Z^2 = -\mu^2 + \frac{m_{H_d}^2(M_Z) - m_{H_u}^2(M_Z) \tan^2 \beta}{\tan^2 \beta - 1}. \quad (1)$$

where $\tan \beta = v_u/v_d$ is the ratio of vacuum expectation values of H_u and H_d . For $\tan \beta$ larger than $\gtrsim 5$ the formula simplifies to

$$\frac{1}{2}M_Z^2 \simeq -\mu^2(M_Z) - m_{H_u}^2(M_Z). \quad (2)$$

The EW scale value of $m_{H_u}^2$ depends on the boundary condition at a high scale and the correction accumulated through the renormalization group (RG) evolution:

$$m_{H_u}^2(M_Z) = m_{H_u}^2 + \delta m_{H_u}^2, \quad (3)$$

and thus we find

$$\frac{1}{2}M_Z^2 \simeq -\mu^2(M_Z) - m_{H_u}^2 - \delta m_{H_u}^2. \quad (4)$$

The RG evolution of μ is controlled only by gauge and Yukawa couplings while the $\delta m_{H_u}^2$ depends on all soft-SUSY-breaking parameters. For given $\tan \beta$ (which sets values of Yukawa couplings), we can solve the RG equations and express EW scale values of $m_{H_u}^2$, μ^2 , and consequently M_Z^2 given by Eq. (2), in terms of all GUT-scale parameters [9, 10] (we consider the GUT-scale as an example, the conclusions do not depend on this choice). For $\tan \beta = 10$, we have:

$$M_Z^2 \simeq -1.9\mu^2 + 5.9M_3^2 - 1.2m_{H_u}^2 + 1.5m_t^2 - 0.8A_tM_3 + 0.2A_t^2 + \dots, \quad (5)$$

where parameters appearing on the right-hand side are the GUT-scale parameters (we do not write the scale explicitly). Here, M_3 is the $SU(3)$ gaugino mass, A_t is the top soft-SUSY-breaking trilinear coupling, and for simplicity we have defined $m_{\tilde{t}}^2 \equiv (m_{\tilde{t}_L}^2 + m_{\tilde{t}_R}^2)/2$ and assumed that the soft-SUSY-breaking stop mass-squared parameters are comparable, $m_{\tilde{t}_L}^2 \simeq m_{\tilde{t}_R}^2$. Other scalar masses and the $SU(1)$ and $SU(2)$ gaugino masses, M_1 and M_2 , appear with negligible coefficients and we neglect them in our discussion. The coefficients in this expression depend only on $\tan\beta$ (they do not change dramatically when varying $\tan\beta$ between 5 and 50) and $\log(M_{\text{GUT}}/M_Z)$. In a similar way we can express the EW-scale values of the stop mass-squared, gluino mass and top trilinear coupling in terms of GUT-scale boundary conditions. For $\tan\beta = 10$ we have:

$$m_{\tilde{t}}^2(M_Z) \simeq 5.0M_3^2 + 0.6m_{\tilde{t}}^2 + 0.2A_tM_3 \quad (6)$$

$$M_3(M_Z) \simeq 3M_3 \quad (7)$$

$$A_t(M_Z) \simeq -2.3M_3 + 0.2A_t. \quad (8)$$

Comparing Eq. (5) with Eqs. (6) and (7) we easily see the usual expectation from SUSY:

$$M_Z \simeq m_{\tilde{t}_{1,2}} \simeq m_{\tilde{g}} \simeq \mu, \quad (9)$$

when all the soft-SUSY-breaking parameters are comparable. Furthermore, neglecting terms proportional to A_t in Eqs. (8) and (6) we find that the *typical stop mixing* is

$$\frac{|A_t|}{m_{\tilde{t}}}(M_Z) \simeq \frac{2.3M_3}{\sqrt{5.0M_3^2 + 0.6m_{\tilde{t}}^2}} \lesssim 1.0, \quad (10)$$

which has an important implication for the Higgs mass.

B. Typical Higgs mass

The minimal supersymmetric model (MSSM) contains two Higgs doublets which result in five Higgs bosons in the spectrum: light and heavy CP even Higgses, h and H , the CP odd Higgs, A , and a pair of charged Higgs bosons, H^\pm . In the decoupling limit, $m_A \gg m_Z$, we find that $m_A \simeq m_H \simeq m_{H^\pm}$ and the light CP even Higgs has the same coupling to the Z-boson as the Higgs in the standard model; we say the light CP even Higgs boson is standard-model-like (SM-like). Its mass is approximately given as:

$$m_h^2 \simeq M_Z^2 \cos^2 2\beta + \frac{3G_F m_t^4}{\sqrt{2}\pi^2} \left\{ \log \frac{m_{\tilde{t}}^2}{m_t^2} + \frac{A_t^2}{m_{\tilde{t}}^2} \left(1 - \frac{A_t^2}{12m_{\tilde{t}}^2} \right) \right\}. \quad (11)$$

where the first term is the tree level result and the second term is the dominant one-loop correction [11, 12, 13, 14].

In the SM the dominant decay mode of the Higgs boson is $h_{\text{SM}} \rightarrow b\bar{b}$ when $m_{h_{\text{SM}}} \lesssim 140$ GeV. This is also the dominant decay mode of the SM-like Higgs in most of the SUSY parameter space in the MSSM. At LEP the SM-like Higgs boson could be produced in association with the Z-boson and LEP has placed strong constraints on $Zh \rightarrow Zb\bar{b}$. The limits on

$$C_{eff}^{2b} \equiv [g_{ZZh}^2/g_{ZZh_{\text{SM}}}^2]\text{Br}(h \rightarrow b\bar{b}), \quad (12)$$

where g_{ZZh} is the $Z-Z-h$ coupling in a given model, are shown in Fig. 1 (from Ref. [7, 15]). From this plot, one concludes that $m_h < 114$ GeV is excluded for a SM-like h that decays primarily to $b\bar{b}$.²

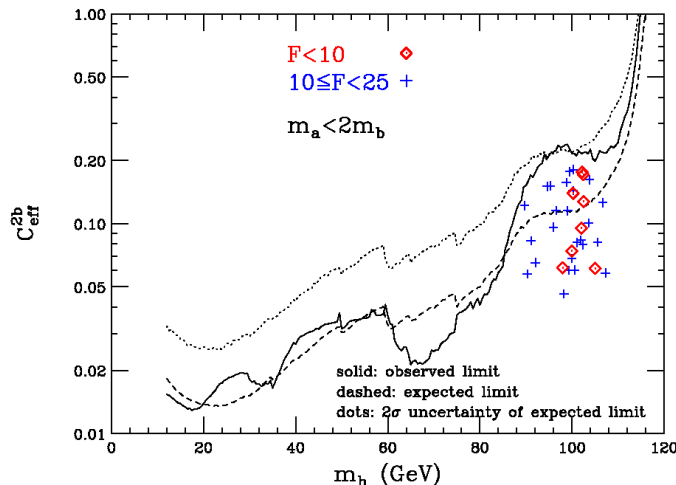


FIG. 1: Expected and observed 95% CL limits on C_{eff}^{2b} are shown vs. m_h . Also plotted are the predictions for the NMSSM parameter cases having fixed $\tan\beta = 10$, $M_{1,2,3}(M_Z) = 100, 200, 300$ GeV that give fine-tuning measure $F < 25$ and $m_{a_1} < 2m_b$ and that are consistent with Higgs constraints.

From, Eq. (11) we see that at tree level, $m_h < M_Z \simeq 91$ GeV which does not satisfy the LEP limit. The Higgs mass can be increased beyond this value either by increasing the

² For discussion of the possibility that H is SM-like or that h and H share the coupling to ZZ and WW see *e.g.* [16] and references therein. This possibility does not have a significant impact on the discussion of naturalness of EWSB which is the focus of this review.

mixing in the stop sector, $A_t/m_{\tilde{t}}$, or by increasing the stop mass, $m_{\tilde{t}}^2$. The effect of stop mixing on the Higgs mass is shown in Fig. 2 obtained from *FeynHiggs-2.5.1* [17, 18]. As we have learned, the typical mixing in the stop sector, which is achieved as a result of RG evolution from a large range of high scale boundary conditions, is $|A_t|/m_{\tilde{t}}(M_Z) \lesssim 1.0$. With this typical mixing and superpartners near the EW scale, we obtain the *typical Higgs mass*, $m_h \simeq 100$ GeV, see Fig. 3. This is a prediction from a large range of SUSY parameter space and it is highly insensitive to small variations of soft SUSY breaking terms at the unification scale (for fixed mixing the Higgs mass depends only logarithmically on stop masses).

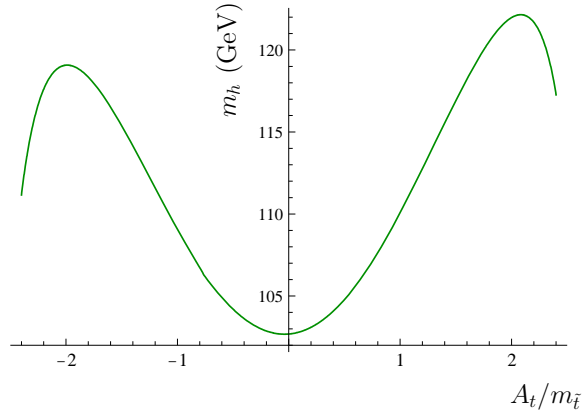


FIG. 2: Mass of the Higgs boson in the MSSM as a function of the mixing in the stop sector, $A_t/m_{\tilde{t}}$, for $\tan\beta = 10$ and $m_{\tilde{t}} = 400$ GeV ($M_{SUSY} = m_A = \mu = 400$ GeV) obtained from *FeynHiggs-2.5.1*.

In order to push the Higgs mass above the LEP limit, 114.4 GeV, assuming the typical mixing, the stop masses have to be $\gtrsim 1$ TeV.³

Before we connect the discussion of naturalness of EWSB and the Higgs mass it is worthwhile to note an indirect prediction for the SM Higgs mass from precision electroweak data. The latest best fit is obtained for $m_h = 90$ GeV with an experimental uncertainty of +36 and -27 GeV (at 68 percent confidence level) [20]. Excluding the most discrepant measurement, the forward-backward asymmetry for b quarks – A_{FB}^b (with a pull of about 3σ), the best fit

³ The Higgs mass is maximized for $|A_t|/m_{\tilde{t}} \simeq 2$, which corresponds to the maximal-mixing scenario. In this case, $m_{\tilde{t}}$ can be as small as ~ 300 GeV without violating the bound on m_h from LEP. However it is not trivial to achieve the maximal-mixing scenario in models. For more details see *e.g.* the discussion in Refs. [15, 19].

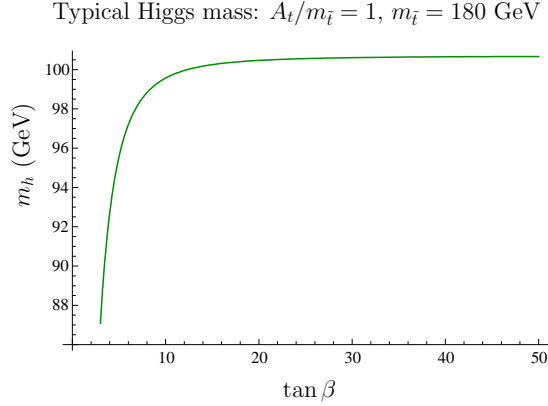


FIG. 3: Mass of the Higgs boson in the MSSM as a function of $\tan \beta$ for typical mixing in the stop sector, $-A_t/m_{\tilde{t}} = 1$, with $M_{SUSY} = m_A = \mu = 200$ GeV obtained from *FeynHiggs-2.5.1*.

value for the Higgs mass is significantly lower which increases the tension with the direct search limit [21].

C. Fine Tuning Problem

From the discussion above we can easily see the tension between the direct search bound on the Higgs mass and naturalness of electroweak symmetry breaking in the MSSM.

In order to push the Higgs mass above the LEP limit, 114.4 GeV, assuming the typical mixing in the stop sector, the stop masses have to be $\gtrsim 1$ TeV. With 1 TeV stop masses one would expect the mass of the Z boson to be of order 1 TeV, see Eq. (9). Or alternatively, with the natural expectation for the stop masses from electroweak symmetry breaking, one would obtain the mass of the Higgs boson $\lesssim 100$ GeV, which is ruled out by LEP limit.

Stop masses of order 1 TeV mean that there is a term of order $(1 \text{ TeV})^2$ on the right hand side of Eq. (5), or that the $\delta m_{H_u}^2$ in Eq. (4) is of order $(1 \text{ TeV})^2$.⁴ Then in order to obtain $M_Z^2 = (91 \text{ GeV})^2$ on the left-hand side of these equations we need to cancel the large

⁴ This result can be obtained by integrating the RG equation for H_u keeping only the term proportional to stop masses, $\delta m_{H_u}^2 \simeq -3\lambda_t^2/(4\pi^2) m_{\tilde{t}}^2 \log(M_{GUT}/M_Z)$. Numerically the loop factor times large log is of order one and assuming the stop masses do not change significantly in the RG evolution we find $\delta m_{H_u}^2 \simeq -m_{\tilde{t}}^2$. This gives the correct answer, however it is somewhat oversimplified because it hides (the most important) effect of the gluino mass on the evolution of stop masses and consequently the Higgs soft mass parameter which is clearly visible in Eqs. (6) and (5) and can be obtained only by solving coupled RG equations.

contribution on the right-hand side resulting from the stop masses by something else with about 1% precision. In Eq. (4) it can be either the boundary condition at the GUT scale or the μ parameter. In both cases small changes in boundary conditions would generate a very different value for the EW scale.

The situation improves when considering large mixing in the stop sector. The mixing is controlled by the ratio of $A_t - \mu \cot \beta$ and $m_{\tilde{t}}$. Since we consider parameter space where μ is small to avoid fine tuning and $\tan \beta \gtrsim 5$ in order to maximize the tree level Higgs mass, see Eq. (11), the mixing is simply given by $A_t/m_{\tilde{t}}$. It was realized that mixing $A_t(M_Z)/m_{\tilde{t}}(M_Z) \simeq \pm 2$ maximizes the Higgs mass for given $m_{\tilde{t}}$ [22], while still satisfying constraints to avoid charge and color breaking (CCB) minima [23]. Using *FeynHiggs-2.2.10* we find that $m_{\tilde{t}}(M_Z) \simeq 300$ GeV and $|A_t(M_Z)| = 450$ GeV (for $\tan \beta \gtrsim 50$), $|A_t(M_Z)| = 500$ (for any $\tan \beta \gtrsim 8$) or $|A_t(M_Z)| = 600$ GeV (for $\tan \beta$ as small as 6) satisfies the LEP limit on the Higgs mass. Therefore large stop mixing, $|A_t(M_Z)/m_{\tilde{t}}(M_Z)| \gtrsim 1.5$ is crucial for satisfying the LEP limit with light stop masses (the physical stop mass in this case can be as small as current experimental bound, $m_{\tilde{t}_1} \gtrsim 100$ GeV). Decreasing the mixing requires increasing of $m_{\tilde{t}}$ and finally we end up with $m_{\tilde{t}} \gtrsim 1$ TeV for small mixing.

In the limit when the stop mass, $m_{\tilde{t}}(M_Z) \simeq 300$ GeV, originates mainly from M_3 , from Eq. (6) we see we need $M_3 \simeq 130$ GeV. Then Eq. (8) shows that the necessary $|A_t(M_Z)| \simeq 500$ GeV is obtained only when $A_t \lesssim -1000$ GeV or $A_t \gtrsim 4000$ GeV at the GUT scale, in both cases it has to be significantly larger than other SSBs. The contribution from the terms in Eq. (5) containing M_3 and A_t is at least $(600 \text{ GeV})^2$ and therefore large radiative correction have to be cancelled either by μ^2 or $m_{H_u}^2(M_{GUT})$. If $m_{\tilde{t}}$ is not negligible at the GUT scale, M_3 can be smaller, but in this case we need even larger A_t and the conclusion is basically the same. Thus, although the situation improves by considering large A_t term, we still need at least 3 % fine tuning [19].⁵

To summarize, in the MSSM with unrelated soft-susy-breaking terms the LEP limit on the Higgs mass requires cancellations between soft-susy-breaking parameters and/or the μ parameter at the level of $\sim 1\%$ (optimizing the mixing in the stop sector reduces the required cancellations to $\sim 3\%$).

⁵ For additional discussion of the connection between the stop sector and the Higgs boson mass and related issues of fine tuning of EWSB see also Refs. [24, 25].

D. Naturalness as a guide

Naturalness of EWSB in supersymmetric models is the only reason to expect superpartners at the LHC. If naturalness or the fine-tuning criterion are not relevant, models like MSSM can be realized in nature with superpartners beyond the LHC reach, $\gtrsim 3$ TeV, while still providing all the virtues discussed in the introduction, namely gauge coupling unification⁶, radiative EWSB, even the explanation of the observed dark matter density. And all this requires only fine tuning at the level of 0.1%.

The acceptable level of fine tuning is clearly a subjective criterion⁷ and it is certainly not a reason to abandon some models or ideas. After all, the smallness of the electroweak scale might be just an accident without any explanation, or simply the whole idea of low scale supersymmetry is wrong and there is a different mechanism realized in nature that is responsible for EWSB.

For the purposes of this review we will nevertheless assume that the low scale supersymmetry is the scenario responsible for EWSB, and more importantly, the EWSB happens in a natural way without necessity of fine-tuning between radiative corrections and boundary values of relevant parameters. This is the premise of this review and we will focus on the consequences this assumption has for the Higgs sector and/or SUSY breaking scenarios. We will see that the naturalness criterion will lead us to consider models in which the Higgs boson decays in an unusual way or models with unusual pattern of soft supersymmetry breaking parameters.⁸

III. MODELS WITH NON-STANDARD HIGGS DECAYS

As we saw the fine-tuning problem in EWSB in the MSSM doesn't result from non-observation of superpartners (current experimental limits on relevant superpartners, see Eq. (9), are not far from the EW scale) but rather from the non-observation of the Higgs boson.

⁶ Somewhat heavy superpartners typically lead to larger splitting between masses of squarks and sleptons which can even improve the precision with which the gauge couplings unify through the weak scale threshold corrections [26].

⁷ For an interesting review of various views of naturalness see *e.g.* Ref. [27].

⁸ Due to limited space we do not discuss models that focus on increasing the mass of the Higgs boson by additional interactions, *e.g.* gauge interaction, see Refs. [28, 29].

This suggests that the fine-tuning problem could be eliminated if one can avoid the LEP limits on the Higgs mass [30].

The solution to the fine-tuning problem in models in which the SM-like Higgs decays dominantly in a different way is straightforward. If the $h \rightarrow b\bar{b}$ decay mode is not dominant the Higgs boson does not need to be heavier than 114 GeV, it can be as light as the typical Higgs mass or even lighter depending on the limits placed on the dominant decay mode. If this limit is ≤ 100 GeV, there is no need for large superpartner masses and superpartners can be as light as current experimental limits allow [30]. For a review of the experimental limits on the mass of the SM-like Higgs boson in various decay modes see Ref. [31]. Quite surprisingly, only if the Higgs decays primarily to two or four bottom quarks, two jets, two taus or to an invisible channel (such as two stable LSP's), is the LEP limit on m_h above 100 GeV. LEP limits on m_h for all other decay modes are below 90 GeV and would therefore not place a constraint on superpartner masses. Thus, models where these alternate (less constrained) decay modes are dominant provide a solution to the fine tuning problem.⁹

A. NMSSM with a light singlet-like CP odd Higgs

The situation we have just discussed already happens in the simplest extension of the MSSM - the next-to-minimal supersymmetric model (NMSSM). The NMSSM introduces an additional singlet super-field with the following couplings

$$W \supset \lambda \widehat{S} \widehat{H}_u \cdot \widehat{H}_d + \frac{1}{3} \kappa \widehat{S}^3. \quad (13)$$

The vacuum expectation value of the singlet effectively generates the μ term. The Higgs spectrum consist of three CP even Higgses, $h_{1,2,3}$, two CP odd Higgses, $a_{1,2}$ and a pair of charged Higgs bosons. It was found that in the NMSSM the SM-like Higgs boson can be light, $m_{h_1} \sim 100$ GeV, as is predicted from natural EWSB with a generic SUSY spectrum. The Higgs in this scenario decays partly via $h_1 \rightarrow b\bar{b}$ but dominantly into two CP-odd Higgs bosons, $h_1 \rightarrow a_1 a_1$, where $m_{a_1} < 2m_b$ so that $a_1 \rightarrow \tau^+ \tau^-$ (or light quarks and gluons) decays are dominant [15, 30]. Further, since the ZZh and $WW h$ couplings are very SM-like, these

⁹ In specific models, avoiding the fine tuning problem might require another tuning of parameters in order to make an alternate decay mode for the Higgs boson dominant, see e.g. Refs. [32, 33] for the discussion of these issues in the NMSSM.

scenarios give excellent agreement with precision electroweak data.

Besides alleviating or completely removing the fine tuning problem the possibility of modified Higgs decays is independently supported experimentally. The largest excess (2.3σ) of Higgs-like events at LEP is that in the $b\bar{b}$ final state for reconstructed $M_{b\bar{b}} \sim 98$ GeV. The number of excess events is roughly 10% of the number of events expected from the standard model with a 98 GeV Higgs boson. Thus, this excess cannot be interpreted as the Higgs of the standard model or the SM-like Higgs of the MSSM.¹⁰ However, this excess is a perfect match to the idea of non-standard Higgs decays. As we have discussed, from natural EWSB we expect the SM-like h to have mass very near 100 GeV, and this is possible in any model where the SM-like Higgs boson decays mainly in a mode for which the LEP limits on m_h are below 90 GeV. The $h \rightarrow b\bar{b}$ decay mode will still be present, but with reduced branching ratio. Any $\text{Br}(h \rightarrow b\bar{b}) \lesssim 30\%$ is consistent with experimental limits for $m_h \sim 100$ GeV. Further, $\text{Br}(h \rightarrow b\bar{b}) \sim 10\%$ with $m_h \sim 100$ GeV provides a perfect explanation of the excess. This interpretation of the excess was first made in the NMSSM scenario discussed above with the $h_1 \rightarrow a_1 a_1 \rightarrow \tau^+ \tau^- \tau^+ \tau^-$ mode being dominant [34], but it clearly applies to a wide variety of models. The NMSSM scenarios with minimal fine-tuning, and thus with the SM-like Higgs mass at ~ 100 GeV, nicely explain the excess, see Fig. 1 (from Ref. [15]).

In summary, models with non-standard Higgs decays can avoid the fine-tuning problem by allowing the Higgs boson mass to be ~ 100 GeV, value predicted from natural EWSB, while at the same time the subdominant decay mode, $h \rightarrow b\bar{b}$, with $\sim 10\%$ branching ratio can explain the largest excess of Higgs-like events at LEP at $M_{b\bar{b}} \sim 98$ GeV. A SM-like h with $m_h \sim 100$ GeV is also nicely consistent with precision electroweak data. Besides the NMSSM scenario discussed above, non-standard Higgs decays occur in a variety of models, for a review see Ref. [31]. An interesting model featuring all the virtues of the above scenario with the SM-like Higgs decaying into four gluons appeared recently [35].

¹⁰ In the MSSM this excess can be explained by the Higgs with highly reduced coupling to ZZ . This explanation doesn't remove the fine tuning problem since it is the heavy Higgs which is SM-like and has to satisfy the 114 GeV limit. For a detailed discussion and references, see Ref. [16].

B. Light doublet-like CP odd Higgs at small $\tan\beta$

An interesting variation of the above NMSSM scenario is the scenario with a *doublet-like* CP odd Higgs below the $b\bar{b}$ threshold. For small $\tan\beta$, $\tan\beta \lesssim 2.5$, it is the least constrained (and only marginally ruled out) in the MSSM, and thus easily viable in simple extensions of the MSSM [36]. Surprisingly the prediction from this region is that all the Higgses resulting from two Higgs doublets: h , H , A and H^\pm could have been produced already at LEP or the Tevatron, but would have escaped detection because they decay in modes that have not been searched for or the experiments are not sensitive to.

The heavy CP even and the CP odd Higgses could have been produced at LEP in $e^+e^- \rightarrow HA$ but they would avoid detection because H dominantly decays to ZA - the mode that has not been searched for. The charged Higgs is also very little constrained and up to $\sim 40\%$ of top quarks produced at the Tevatron could have decayed into charged Higgs and the b quark since the dominant decay mode for the charged Higgs $H^\pm \rightarrow W^\pm A$ with $A \rightarrow c\bar{c}$ or $\tau^+\tau^-$ was not searched for either. In addition the charged Higgs with the properties emerging in this scenario and mass close to the mass of the W boson could explain the 2.8σ deviation from lepton universality in W decays measured at LEP [37] as discussed in Ref. [38].

The mass of the light CP even Higgs is the only problematic part in this scenario. In the MSSM in this region of the parameter space we find $m_h \simeq 40 - 60$ GeV, and thus $h \rightarrow AA$ decay mode is open and generically dominant. However the decay mode independent limit requires the SM like Higgs to be above 82 GeV which rules this scenario out in the MSSM, since m_h cannot be pushed above 82 GeV by radiative corrections. There are however various ways to increase the mass of the SM-like Higgs boson in extensions of the MSSM. A simple possibility is to consider singlet extensions of the MSSM containing $\lambda SH_u H_d$ term in the superpotential. It is known that this term itself contributes $\lambda^2 v^2 \sin^2 2\beta$, where $v = 174$ GeV, to the mass squared of the CP even Higgs [39] and thus can easily push the Higgs mass above the decay-mode independent limit, 82 GeV. Note, this contribution is maximized for $\tan\beta \simeq 1$.¹¹ Thus it is not surprising that the scenario with a light doublet-like CP-odd

¹¹ Sometimes it is argued that the extra contribution to the Higgs mass in the NMSSM makes the NMSSM less fine-tuned than the MSSM since the LEP limit 114 GeV is easier to satisfy. This effect is not very significant however. For large $\tan\beta$ the extra contribution is negligible and for small $\tan\beta$ we are losing the tree level contribution we have in the MSSM, see Eq. (11). Optimizing $\tan\beta$ improves the naturalness of EWSB in the NMSSM compared to the MSSM only marginally [15].

Higgs boson is phenomenologically viable in the NMSSM [40] but clearly since the light CP-odd Higgs comes from the two Higgs doublets it is not limited to singlet extensions of the MSSM.

C. Higgs as a Link to New Sectors

The extra singlet in the NMSSM does not spoil any of the virtues of the MSSM, including the possibility of gauge coupling unification and matter particles fitting into complete GUT multiplets. It is also very interesting to consider a possibility that besides the extra singlet there is a whole extra (E) sector of particles (for simplicity we assume that these particles are scalars) that are singlets under standard model gauge symmetry and couple to the MSSM sector only through the Higgs fields. Such couplings would have a negligible effect on the phenomenology involving SM matter particles, whereas they can dramatically alter Higgs physics. For example, they would allow the lightest CP-even Higgs boson h to decay into two of the E -particles if the E -particles are light enough.

If the h decays to two lighter E -particles, then the strategy for Higgs discovery will depend on the way the E -particles appearing in the decays of the h themselves decay. They might decay predominantly into other stable E -particles, in which case the MSSM-like h decays mainly invisibly. More typically, however, the E -particles acquire a tiny coupling to SM particles via interactions with the MSSM Higgs fields (scalar mass eigenstates are generically mixed). In this case, light E -particles will decay into $b\bar{b}$, $\tau^+\tau^-$ or other quarks or leptons depending on the model. Although E -particles would have small direct production cross sections and it would be difficult to detect them directly, their presence would be manifest through the dominant Higgs decay modes being $h \rightarrow 4f$, where $4f$ symbolically means four SM particles, *e.g.* $b\bar{b}b\bar{b}$, $b\bar{b}\tau^+\tau^-$, $\tau^+\tau^-\tau^+\tau^-$, 4γ and so on. The situation can be even more complicated if the h decays to E -particles that themselves decay into other E -particles, \dots , which in turn finally decay into SM particles. In this case the SM-like Higgs boson would cascade decay in the E -sector till it decays to the lightest E -particles, which finally decay to SM particles. Depending on the mass of the lightest E -particle this would look like the SM-like Higgs is decaying into a large number of light jets, or a large number of muons, or a large number of electrons, or in case the lightest E -particle is below $2m_e$ threshold the only way it can decay is into two photons, so effectively the SM-like Higgs

boson would burst into a large number of soft photons [42]. These events would be quite spectacular, e.g. the decay of the Higgs boson to many photons would light up a large portion of the electromagnetic calorimeter. Such signatures are not very usual in particle physics which is probably the reason these were never looked for.

The light CP-odd Higgs of the NMSSM or other extended models might be within the reach of current B factories where it can be produced in Upsilon decays, $\Upsilon \rightarrow A\gamma$ [41]. The same applies to light E -sector particles¹² with all the possible spectacular decay modes mentioned above [42].

IV. UNUSUAL SUPERSYMMETRY

Another possibility to reconcile non-observation of the Higgs boson with naturalness of electroweak symmetry breaking assumes that the MSSM is the correct description of nature at the TeV scale but the usual assumptions we make about SUSY breaking are not correct. In this case in order to radiatively trigger EWSB and satisfy the limit on the Higgs mass SUSY has to be quite special.

Although M_Z results from cancellations between SSB parameters in this case, it does not necessarily mean that the Z mass is fine-tuned. SUSY breaking scenarios typically produce SSBs which are related to each other in a specific way and so they should not be treated as independent parameters. In such a case, the discussion around Eq. (5) is not relevant to assess the level of fine tuning, however it still can be used to explore relations between SSBs that have to be generated, should the M_Z emerge in a natural way.

A simple example of a relation between SSBs that allows relevant SSBs to be large without significantly contributing to M_Z is $m_{\tilde{t}} = m_{H_u}$. In this case the contribution from $m_{\tilde{t}}$ approximately cancels the contribution from m_{H_u} in Eq. (5) and this cancelation is more precise for large $\tan\beta$ as discussed in Ref. [43]. Thus scalars can be much heavier than gauginos without contributing more to M_Z .

Even larger hierarchy between scalars and gauginos can be achieved with more optimized relations between SSBs: $A_0 \simeq -2m_{16}$ and $m_{10} \simeq \sqrt{2}m_{10}$ with $M_{1/2} \ll m_{16}$, where A_0 is the universal soft trilinear coupling at the GUT scale, and m_{16} and m_{10} are universal

¹² The light CP-odd Higgs of the NMSSM is an example of an E -sector with only one particle.

masses of superpartners of three families and the Higgs masses respectively. These conditions, motivated by SO(10) grand unified theories, were found to radiatively generate the largest hierarchy between the first two generations of scalars and the third one [44] and were also considered for having attractive features, including the possibility of Yukawa coupling unification [45, 46] and maximizing the suppression of the proton decay [47]. However it should be stressed that these type of relations could reduce fine-tuning of EWSB only if they can be realized as a unique (or at least a discrete) possibility in a given SUSY breaking scenario, which proved not to be easy.

More recent example is a proper mixture of anomaly and modulus mediation [48, 49, 50, 51] which produces boundary conditions such that the initial value of $m_{H_u}^2$ is canceling most of the contribution from running.

All the above mentioned examples have one thing in common: the boundary conditions for soft SUSY breaking parameters generated by some mechanism are such that the boundary condition for $m_{H_u}^2$ is canceling the accumulated radiative correction $\delta m_{H_u}^2$ from the RG evolution, see Eq. (4). Clearly such solutions must exist, and some of them might be even realized in simple models but if this cancellation is required to be very precise then there is a “coincidence” problem: the relations that have to be satisfied between SSBs in order to recover the correct M_Z depend on the energy interval SSBs are going to be evolved over. Therefore a SUSY breaking scenario would have to know that SSBs will evolve according to MSSM RG equations, and exactly from e.g. M_{GUT} to M_Z . For example, the top Yukawa coupling is responsible for driving the $m_{H_u}^2$ to negative values and $\delta m_{H_u}^2$ directly depends on its value. In models we envisioned so far the generated boundary condition for $m_{H_u}^2$ does not depend on the value of the top Yukawa coupling and thus it would be a “coincidence” that after RG evolution $\delta m_{H_u}^2$ precisely cancels the boundary condition at a high scale. In other words the M_Z in this case is highly sensitive to the value of the top Yukawa coupling.¹³

In this section we will discuss a possibility that $\delta m_{H_u}^2 \simeq O(M_Z)$. Therefore none of the terms in Eq. (4) are large and no cancellation, or fine-tuning, is necessary to reproduce the correct value of M_Z . The cancellations (a relations between SSBs supplied by a given model) might still be required to achieve $\delta m_{H_u}^2 \simeq O(M_Z)$. Nevertheless these scenarios could avoid the coincidence problem mentioned above. A typical feature of these scenarios is the

¹³ Similar discussion applies to the sensitivity of M_Z to α_3 .

reduced sensitivity of M_Z to dimensionless couplings, especially the top Yukawa coupling. Additional interesting feature is that a large hierarchy between the fundamental scale at which SSB are generated, e.g. the GUT scale or the Planck scale, and the EW scale is required in order for these scenarios to be viable.

A. Negative stop masses squared

A simple way to achieve $\delta m_{H_u}^2 \simeq O(M_Z)$ is to consider negative stop masses squared at a high scale [19]. If we allow negative stop masses squared at the GUT scale several interesting things happen simultaneously. First of all, from Eq. (6) we see that unless $m_{\tilde{t}}$ is too large compared to M_3 it will run to positive values at the EW scale. At the same time the contribution to $m_{H_u}^2$ from the energy interval where $m_{\tilde{t}}^2 < 0$ partially or even exactly cancels the contribution from the energy interval where $m_{\tilde{t}}^2 > 0$ and so the EW scale value of $m_{H_u}^2$ can be arbitrarily close to the starting value at M_{GUT} , see Fig. 4. From Eq. (5) we see that this happens for $m_{\tilde{t}}^2 \simeq -4M_3^2$ (neglecting A_t). No cancellation between initial value of $m_{H_u}^2$ (or μ) and the contribution from the running is required. And finally, from Eqs. (6) and (8) we see that the stop mixing is typically much larger than in the case with positive stop masses squared. For positive (negative) stop masses squared we find $|A_t(M_Z)/m_{\tilde{t}}(M_Z)| \lesssim 1$ ($\gtrsim 1$) starting with $A_t = 0$ and small $m_{\tilde{t}}$ at the GUT scale. Starting with larger $m_{\tilde{t}}$ the mixing is even smaller (larger) in the positive (negative) case. Therefore large stop mixing at the EW scale is generic in this scenario and actually it would require very large GUT scale values of A_t to end up with small mixing at the EW scale.

It turns out that in the region where $m_{H_u}^2$ gets negligible contribution from running, the radiatively generated stop mixing is close to maximal even when starting with negligible mixing at the GUT scale. In this case, comparing Eqs. (6) and (8), we find ¹⁴

$$A_t(M_Z)/m_{\tilde{t}}(M_Z) \simeq -1.5 + 0.2A_t/M_3. \quad (14)$$

Slightly more negative stop masses squared at the GUT scale would result in maximal stop mixing at the EW scale even when starting with negligible A_t . Nevertheless the example

¹⁴ To be more precise the generated mixing is somewhat larger than that shown in this equation, since we should minimize the potential at the SUSY scale $\sim m_{\tilde{t}}$ and should not run SSBs all the way to M_Z (see Fig. 4).

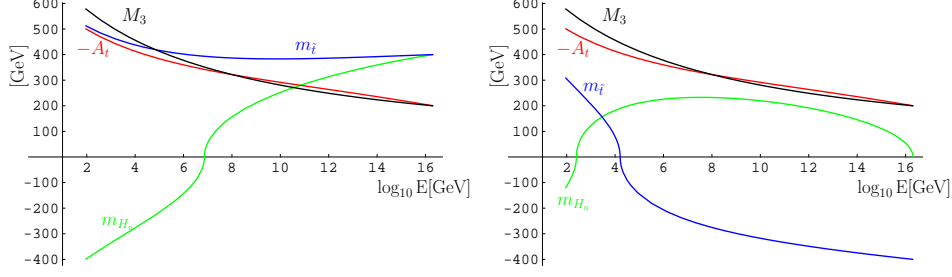


FIG. 4: Right: renormalization group running of relevant SSBs for $\tan\beta = 10$ and GUT scale boundary conditions: $-A_t = M_3 = 200$ GeV, $m_t^2 = -(400 \text{ GeV})^2$ and $m_{H_u}^2 = 0 \text{ GeV}^2$. Left: the same for $m_t^2 = +(400 \text{ GeV})^2$; the $m_{H_u}^2$ is chosen to be the same as m_t^2 in this case in order to keep the same y-axis scale in both plots. In order to have both mass dimension one and two parameters on the same plot and keep information about signs, we define $m_{H_u} \equiv m_{H_u}^2 / \sqrt{|m_{H_u}^2|}$ and $m_t \equiv m_t^2 / \sqrt{|m_t^2|}$.

in Fig. 4 with simple GUT scale boundary conditions already leads to EW scale parameters $m_t(M_Z) \simeq 300$ GeV and $A_t(M_Z) = -500$ GeV producing sufficiently heavy Higgs boson, $m_h \simeq 115.4$ GeV. Small variations of GUT scale parameters, including positive or negative values of $m_{H_u}^2$, would produce similar results and scaling all parameters up would lead to larger Higgs mass.

In a theory which predicts $m_t^2 \simeq -4M_3^2$, the fine tuning problem is entirely solved. The contribution to $m_{H_u}^2$ from the running is negligible and the $\mathcal{O}(M_Z^2)$ values of $m_{H_u}^2$ and μ^2 at the GUT scale naturally result in the correct M_Z . However, the absence of fine tuning is quite robust and the relation above does not have to be satisfied very precisely. Requiring fine tuning less than 10% the stop masses squared have to be generated between $\simeq -3M_3^2$ and $\simeq -5M_3^2$ for $M_3 \simeq 200$ GeV. This interval is shrinking with increasing M_3 which is a sign of the coincidence problem discussed above.

Models with negative stop masses at a high scale also reduce the sensitivity of the EW scale to the top Yukawa coupling. This can be qualitatively understood from the fact that it is the top Yukawa coupling that drives the $m_{H_u}^2$ to positive values in the energy interval where stop masses are negative, and it is the same top Yukawa coupling that drives it to negative values in the energy interval where stop masses are positive. Thus larger top Yukawa coupling would drive the $m_{H_u}^2$ to higher positive values but then, when stop masses turn positive, it would drive it to negative values faster.

Negative stop masses at a high scale can be realized e.g. in models with gauge messengers [52].

1. Gauge Messenger Model

In the simple gauge messenger model, based on SU(5) supersymmetric GUT with a minimal particle content, it is assumed that an adjoint chiral superfield, $\hat{\Sigma}$, gets a vacuum expectation value in both its scalar and auxiliary components: $\langle \hat{\Sigma} \rangle = (\Sigma + \theta^2 F_\Sigma) \times \text{diag}(2, 2, 2, -3, -3)$. The vev in the scalar component, $\Sigma \simeq M_G$, gives supersymmetric masses to X and Y gauge bosons and gauginos and thus breaks SU(5) down to the standard model gauge symmetry. The vev in the F component, F_Σ , splits masses of heavy gauge bosons and gauginos and breaks supersymmetry. The SUSY breaking is communicated to MSSM scalars and gauginos through loops involving these heavy gauge bosons and gauginos which play the role of messengers (the messenger scale is the GUT scale). The gauge messenger model is very economical, all gaugino and scalar masses are given by one parameter,

$$M_{\text{SUSY}} = \frac{\alpha_G}{4\pi} \frac{|F_\Sigma|}{M_G}, \quad (15)$$

and it is phenomenologically viable [52]. The characteristic features are: negative and non-universal squark and slepton masses squared at the unification scale, non-universal gaugino masses, and sizable soft-trilinear couplings. The combination of a large negative top soft trilinear coupling and negative stop masses square lead to close to maximal mixing scenario for the Higgs mass and reduce the fine-tuning of electroweak symmetry breaking. In this scenario, all soft supersymmetry breaking parameters at the unification scale can be smaller than 400 GeV and all the superpartners can be lighter than 400 GeV and still satisfy all the limits from direct searches for superpartners and also the limit on the Higgs mass.

For soft SUSY breaking parameters in extended models see Ref. [52] and the discussion of dark matter can be found in Ref. [53].

B. Hypercharge Mediation of SUSY Breaking

Another scenario with $\delta m_{H_u}^2 \simeq O(M_Z)$ is the hypercharge mediation of SUSY breaking [54]. It is a string motivated scenario which uses a similar setup envisioned for the anomaly mediation [55, 56] and it is also related to the Z'-mediation [57].

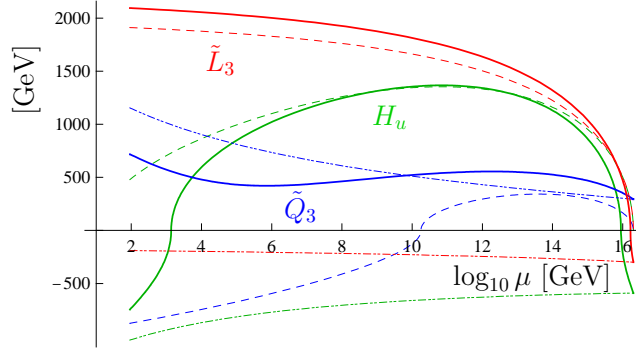


FIG. 5: Renormalization group running of m_{H_u} (green), m_{Q_3} (blue) and m_{L_3} (red) for $\tan \beta = 10$, $m_{3/2} = 50$ TeV and $\alpha = 0.2$ for $M_* = M_{GUT}$. We define $m_{H_u} \equiv m_{H_u}^2 / \sqrt{|m_{H_u}^2|}$ and similarly for m_{Q_3} and m_{L_3} . The contribution of pure hypercharge mediation is given by dashed lines and the separate contribution from anomaly mediation is represented by the corresponding dotted lines.

The scenario is characterized by the bino mass M_1 being the only soft-susy breaking term generated at a high-scale. In the RG evolution to the weak scale, all scalar masses (including the Higgs masses) receive a contribution from the bino mass. This positive contribution dominates at the beginning of the RG evolution. Once sizable scalar masses are developed, the negative contribution from Yukawa couplings becomes important and can overcome the contribution from the bino mass. In pure hypercharge mediation, the left-handed stop mass squared would be driven to negative values, because out of all scalars its hypercharge is the smallest and its Yukawa coupling is the largest. All other squarks and sleptons remain positive, see Fig 5 (from Ref. [54]).¹⁵ The wino and gluino masses receive a contribution from the bino mass at the two loop level.

Due to negative stop masses at the EW scale the scenario is not viable but small contribution to gaugino masses $O(100 \text{ GeV})$ from other sources can easily cure this problem. For example one can consider a subleading contribution from anomaly mediation [54], see Fig 5 (from Ref. [54]). Or, one can consider small contribution to gaugino masses from gravity mediation. The gluino mass large enough to satisfy experimental limits is also sufficient to push stops above the experimental limit and drive the $m_{H_u}^2$ to negative values.

When discussing Eq. (5) we mentioned that the bino mass enters this formula with a

¹⁵ Depending on $\tan \beta$ and the scale at which the SUSY breaking is communicated to the MSSM sector also the up-type Higgs mass squared, $m_{H_u}^2$ can be driven to negative values. However, unless the starting scale is very large (above the GUT scale) $m_{H_u}^2$ remains small and positive at the weak scale.

negligible coefficient and we didn't bother to write it down. However, in the hypercharge mediation it is the only nonzero term (up to possible contribution from other sources). The value of this term is $-0.006 M_1^2$ and thus M_1 of order 4 TeV would not result in fine-tuning larger than 10%. Thus if the pure hypercharge mediation was a phenomenologically viable model, it would be an example of a model in which there is no cancellation required in EWSB. Considering contribution from *e.g.* gravity mediation to the other two gaugino masses of order 200 GeV makes the model viable and does not introduce fine tuning larger than $\sim 4\%$.¹⁶

V. INTERESTING COINCIDENCES?

The LHC will start running soon and we will finally learn about the mechanism responsible for electroweak symmetry breaking. Based on limited available experimental evidence we have envisioned elegant models that are so ambitious that they could describe Nature even at energies close to the Planck scale. These ideas are based on hints or *interesting coincidences* in values of parameters we observe. The most influential such a coincidence is the fact that gauge couplings unify in the MSSM. In addition, the Higgs mass squared parameter, $m_{H_u}^2$, is typically the only one driven to negative values at the EW scale and so SUSY breaking generically results in electroweak symmetry breaking. These can be just coincidences or Nature gave us a hint of what we will really find at the LHC.

Based on similar hints (which are not so clear however) we are guessing how the Higgs sector should look like and how supersymmetry should be broken. It is certainly at least an interesting coincidence that the natural value of the Higgs mass in the MSSM and simple extensions coincides with the largest excess of Higgs like events at LEP. It can be just a coincidence or maybe it is a hint that the Higgs really is at ~ 100 GeV.

Another coincidence is that in models with negative stop masses squared at a high scale the least fine tuned scenarios automatically maximize the Higgs mass through the mixing in the stop sector. In addition, these models, or in general models with $\delta m_{H_u}^2 \simeq O(M_Z)$, including the hypercharge mediation, require a large hierarchy (more than 10 orders of magnitude) between the scale at which SUSY is broken (GUT scale or Planck scale) and

¹⁶ Heavier bino compared to other gauginos proved to be advantageous in earlier studies of non-universal gaugino mediation [58, 59]. However there was no theoretical motivation for such a choice.

the EW scale in order to trigger EWSB.

All these could be just coincidences and some of them certainly are. We put a lot of effort to build a beautiful picture out of what we know. If none of this is realized in Nature, and we were completely misled, let us hope that Nature prepared for us something even more elegant.

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